Deep Inelastic Scattering from Polarized Deuterons†

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Abstract

The spin-dependent structure function of the deuteron, g_{1D}, is calculated within a covariant framework. The off-shell structure of the bound nucleon gives corrections to the convolution model at a level of half a percent for x below 0.7, increasing to more than five percent at larger x. Overall, the dominant source of error comes from the lack of knowledge associated with the deuteron D-state, which may introduce an uncertainty in the neutron spin structure function, g_{1n}, extracted from deuterium data of up to ten percent for x around 0.2.

ADP-94-16/T157

TPR-94-22

To appear in Physics Letters B.

† Work supported in part by BMFT grant 06 OR 735.

There is currently much discussion about the interpretation of the results from the Spin Muon Collaboration [1] and SLAC E142 [2] experiments on the polarized deuteron and helium structure functions. Of particular interest is the neutron structure function, g_{1n}. A precise knowledge of this is necessary to verify the Bjorken sum rule — when combined with the previously measured proton structure function g_{1p} [3, 4, 5, 6]. In addition, an independent measurement of g_{1n} is important for determining the flavor singlet combination of polarized quark distributions — the first moment of which, in the naive parton model, is just the fraction of the spin of the nucleon carried by quarks.

Before the nuclear data can be used for these purposes, it is essential to account for any nuclear effects that may arise when extracting g_{1N} from g_{1A}. Nuclear corrections to polarized deuteron and helium structure functions, such as those due to Fermi motion, shadowing or meson-exchange currents, have been considered by several authors [7, 8, 9, 10, 11, 12]. Early attempts [8] to describe nuclear effects in the deuteron, based on time-ordered perturbation theory in the infinite momentum frame, suffered from the problem that the deuteron wavefunctions are not known in this frame. Subsequent analyses [9, 10] all adopted the so-called convolution approach, in which the nuclear structure function is a one-dimensional convolution of the structure function of a free nucleon with the nucleon momentum distribution in the nucleus. Efforts to incorporate relativistic effects in the deuteron were made in Refs.[11, 12], however also within the confines of the convolution model.

In contrast with the earlier work, in this letter we demonstrate that in a covariant treatment, inclusion of the full off-mass-shell structure of bound nucleons leads to a breakdown of convolution for spin-dependent nuclear structure functions. A convolution component can, however, be extracted from the full result by selectively taking on-shell limits for the bound nucleon structure function, and neglecting relativistic components of the nuclear wavefunction. Off-shell corrections to the spin-averaged structure function of the deuteron, F_{2D}, were examined in Ref.[13], and found to be about 1-2% for $x < 0.9$ (c.f. heavy nuclei or nuclear matter, where the off-shell effects can be considerably larger [14, 15]). Since the absolute values of g_{1D} and g_{1n} are considerably smaller than F_{2D} or F_{2n}, we may expect nuclear corrections to be of greater relative importance for the spin-dependent structure functions. It is of some importance therefore that the issue of off-shell effects in g_{1D} be seriously addressed.

In the present analysis we restrict ourselves to the valence component of polarized

structure functions, thereby avoiding confronting the issue of the axial anomaly [16], which could divert attention from our main interest. Within the impulse approximation, deep inelastic scattering from a polarized deuteron is described as a two-step process, in terms of the virtual photon–nucleon interaction, parametrized by the truncated antisymmetric nucleon tensor $\widehat{G}_{\mu\nu}(p,q)$, and the polarized nucleon–deuteron scattering amplitude, $\widehat{A}(P,p,S)$. The antisymmetric part of the deuteron hadronic tensor can then be written as:

$$\begin{aligned} M_D\, W^D_{\mu\nu}(P,S,q) \;&\equiv\; i\frac{M_D}{P\cdot q}\epsilon_{\mu\nu\alpha\beta}\, q^\alpha \left(S^\beta\; g_{1D}(P,q) \;+\; \left(S^\beta - \frac{S\cdot q}{P\cdot q}P^\beta \right) g_{2D}(P,q) \right) \\ &=\; \int \frac{d^4p}{(2\pi)^4} \mathrm{Tr}\left[\widehat{A}(P,S,p)\;\; \widehat{G}_{\mu\nu}(p,q) \right]\; 2\pi\delta\left((P-p)^2 - M^2 \right), \qquad (1) \end{aligned}$$

where P, p and q are the deuteron, off-shell nucleon and photon momenta, respectively, and M_D is the deuteron mass[1]. The vector S^α ($S^2=-1, P\cdot S = 0$) is defined in terms of deuteron polarization vectors ε^m_α such that [18] $S^\alpha(m) \equiv -i\; \epsilon^{\alpha\beta\lambda\rho}\; \varepsilon^{m*}_\beta\; \varepsilon^m_\lambda\; P_\rho/M_D$, where $m = 0, \pm 1$ is the spin projection.

In analyzing nucleon off-shell effects, it will be convenient to expand the truncated nucleon tensor $\widehat{G}_{\mu\nu}$ in terms of independent basis tensors:

$$\widehat{G}_{\mu\nu}(p,q) \;=\; i\epsilon_{\mu\nu\alpha\beta}\; q^\alpha \left(p^\beta \left(\not{p}\gamma_5\, G_{(p)} \;+\; \not{q}\gamma_5\, G_{(q)} \right) +\; \gamma^\beta\gamma_5\, G_{(\gamma)} \;+\; \cdots \right), \qquad (2)$$

where the coefficients $G_{(i)}$ are scalar functions of p and q, and the dots $(\ldots)$ represent terms that do not contribute to the g_1 structure function in the Bjorken limit, as well as those that vanish in the massless quark limit (which we use throughout).

The polarized structure function, g_{1D}, can be extracted from $W^D_{\mu\nu}$ by considering the polarization combination (m=+1) – (m=–1). Using the fact that $\left(\varepsilon^+_\alpha\; \varepsilon^{+*}_\beta \;-\; \varepsilon^-_\alpha\; \varepsilon^{-*}_\beta \right) = -i\; \epsilon_{\lambda\rho\alpha\beta}\; P^\lambda S^\rho / M_D$ (since $\varepsilon^{+*}_\alpha = -\varepsilon^-_\alpha$), the deuteron–nucleon amplitude $\widehat{A}$ can be written:

$$\widehat{A} \;=\; -\frac{i}{2M_D}\epsilon_{\lambda\rho\alpha\beta}\; P^\lambda S^\rho\; (\not{p}-M)^{-1}\; \Gamma^\alpha(p)\; (\not{P}-\not{p}-M)\; \overline{\Gamma}^\beta(p)\; (\not{p}-M)^{-1}, \qquad (3)$$

where $\Gamma^\alpha(p)$ is the relativistic deuteron–nucleon vertex function [19]. In the massless quark limit only the pseudovector components of $\widehat{A}$ are relevant: $\widehat{A} \equiv \gamma_5\gamma_\lambda\mathcal{A}^\lambda$. In terms of $\mathcal{A}^\lambda$ and $G_{(i)}$, the g_{1D} structure function (per nucleon) is:

$$\begin{aligned} g_{1D}(x) \;&=\; \frac{P\cdot q}{4\pi^2\; M_D\; S\cdot q} \int dy\; dp^2 \left(\mathcal{A}\cdot q\;\; \left(p\cdot q\; G_{(q)} \;+\; G_{(\gamma)} \right) \right. \\ &\qquad\qquad \left. +\; p\cdot q\; \mathcal{A}\cdot p\; G_{(p)} \right), \qquad (4) \end{aligned}$$

[1] For a discussion of some problems associated with using the impulse approximation for g_{2A} see Ref.[17].

where $x = Q^2/2P \cdot q$ is the Bjorken scaling variable and $y = p \cdot q/P \cdot q$ is the light-cone fraction of the deuteron momentum carried by the struck nucleon.

The analogous hadronic tensor for an on-shell nucleon is obtained by tracing $\widehat{G}_{\mu\nu}$ with the nucleon spin-energy projector:

$$2M\ W^N_{\mu\nu}(p,q) \ = \ \mathrm{Tr}\left[(\not p + M)\ \ (1+\gamma_5\ \not s)\ \ \widehat{G}_{\mu\nu}(p,q)\right], \tag{5}$$

where s is the nucleon polarization vector, and by setting $p^2 = M^2$ and $y = 1$ in the coefficients $G_{(i)}$. The structure function g_{1N} can then be identified as $g_{1N}(x) = \widetilde{g}_{1N}(x, p^2 = M^2)$, where

$$\widetilde{g}_{1N}\left(\frac{x}{y}, p^2\right) \ = \ 2p\cdot q\ \left(p \cdot q\ G_{(q)} + G_{(\gamma)}\right). \tag{6}$$

The presence of the $G_{(p)}$ term in Eq.(4) (which does not appear in g_{1N}) means that simple factorization of the fully relativistic g_{1D} into nuclear ($\mathcal{A}^\lambda$) and nucleon ($G_{(i)}$) parts is not strictly possible. This however would be required for convolution, as assumed in Refs.[9, 10, 11]. Nonetheless, by writing $\mathcal{A}^\lambda$ in terms of relativistic deuteron wavefunctions, as calculated for example in [19] (see also [20]), we can show that the $G_{(p)}$ term is of order $(v/c)^3$ compared with the first term in the integrand in Eq.(4). Indeed, all of the non-factorizable corrections to convolution can be shown to be of higher orders in v/c. This is easily seen by separating the "$+$" component of $\mathcal{A}^\lambda$ into an "on-shell" part, which is proportional to the non-relativistic deuteron wavefunctions (see Eq.(12) below), and an off-shell component, $\mathcal{A}^+ \equiv \mathcal{A}^+_{ON} + \mathcal{A}^+_{OFF}$, where (in the deuteron rest frame):

$$\mathcal{A}^+_{ON} \ = \ 2\pi^2 M_D M \left[u^2 + \left(1 - 3\cos^2\theta\right)\frac{uw}{\sqrt{2}} - \left(1 - \frac{3}{2}\cos^2\theta\right) w^2\right], \tag{7}$$

and

$$\begin{aligned}
\mathcal{A}^+_{OFF} \ = \ 2\pi^2 M_D M & \left[\left(\frac{p_z}{M} - \left(1 - \frac{E_p}{M}\right)\cos^2\theta\right)\left(u - \frac{w}{\sqrt{2}}\right)^2\right. \\
& - \frac{3}{2}\left(\frac{p_z}{M} - \frac{E_p}{M}\cos^2\theta\right) v_t^2 \ + \ \frac{3}{\sqrt{2}}\left(1 - \cos^2\theta\right) v_s v_t \\
& + \sqrt{3}\left(\cos\theta - \frac{|\mathbf{p}|}{2M}\left(1 - \cos^2\theta\right)\right) u v_t \\
& - \sqrt{3}\left(\cos\theta - \frac{|\mathbf{p}|}{M}\left(1 - \cos^2\theta\right)\right) w v_t \\
& \left. + \frac{\sqrt{3}|\mathbf{p}|}{M}\left(1 - \cos^2\theta\right)\left(u - \frac{w}{\sqrt{2}}\right) v_s\right],
\end{aligned} \tag{8}$$

with $p_z = |\mathbf{p}|\cos\theta$, $E_p = \sqrt{M^2+\mathbf{p}^2}$ and $\cos\theta = (yM_D - p_0)/|\mathbf{p}|$. In Eqs.(7) and (8) u and w correspond to the S- and D-state deuteron wavefunctions, while v_t and v_s are the relativistic triplet and singlet P-state contributions. (In our numerical calculation we use the model of Ref.[19] with wavefunctions corresponding to the pseudovector πNN interaction.)

Using Eqs.(6)–(8) we can then decompose g_{1D} into a convolution component plus off-shell corrections:

$$g_{1D}(x) = \int_x^1 \frac{dy}{y} \Delta f(y)\ g_{1N}\left(\frac{x}{y}\right)\ +\ \delta^{(N)} g_{1D}(x)\ +\ \delta^{(\mathcal{A})} g_{1D}(x)\ +\ \delta^{(G)} g_{1D}(x), \tag{9}$$

where [9]

$$\Delta f(y) = \int d^4p\ \left(1+\frac{p_z}{M}\right)\ \Delta S(p)\ \delta\left(y - \frac{p^+}{M_D}\right), \tag{10}$$

can be identified with the difference of probabilities to find a nucleon in the deuteron with momentum fraction y and spin parallel and antiparallel to that of the deuteron. For a deuteron with polarization $m=+1$, $\Delta S(p)$ corresponds to the spectral function:

$$\Delta S(p) = \Psi^\dagger_{+1}(p)\ \widehat{S}_z\ \Psi_{+1}(p)\ \ \delta\left(p_0 - M_D + E_p\right), \tag{11}$$

where $\Psi_m(p)$ is the usual (normalized) non-relativistic deuteron wavefunction, and $\widehat{S}_z$ is the z component of the nucleon spin operator. In terms of $\Psi_m(p)$, $\mathcal{A}^+_{ON}$ can be written:

$$\mathcal{A}^+_{ON} = 8\pi^3 M_D M\ \mathcal{N}\ \ \Psi^\dagger_{+1}(p)\ \widehat{S}_z\ \Psi_{+1}(p), \tag{12}$$

where $\mathcal{N} = \int d|\mathbf{p}|\ \mathbf{p}^2\ (u^2+w^2)$ ensures that the normalization agrees with that of the relativistic calculation. The function $\Delta f(y)$ then satisfies $\int_0^1 dy\ \Delta f(y) = 1 - 3/2\ \omega_D$, where $\omega_D = \int d|\mathbf{p}|\ \mathbf{p}^2\ w^2\ /\ \mathcal{N}$ is the non-relativistic D-state probability.

We should point out that the definition of the convolution component in Eq.(9) is not unique. For example, some authors take the non-relativistic limit in the argument of the energy-conserving δ-function, $\delta\left(p_0 - M - \varepsilon_D + \mathbf{p}^2/2M\right)$, where $\varepsilon_D \equiv M_D - 2M$ is the deuteron binding energy. Furthermore, in order to make a comparison of our full result with previous calculations [9], we have included in Eq.(10) the “flux factor” $(1+p_z/M)$, which was introduced in Ref.[9] by analogy with unpolarized structure functions. However, this prescription also differs according to various authors [8, 10, 12]. On the other hand, the full, relativistic result in Eq.(4) is exact (within the impulse approximation) and contains no such ambiguities.

The three non-factorizable corrections in Eq.(9) are proportional to powers of $|\mathbf{p}|/M$. The first two corrections,

$$
\begin{aligned}
\delta^{(\mathcal{A})} g_{1D}(x) \;=\; & \frac{1}{8\pi^2 M_D} \int \frac{dy}{y} dp^2 \left[\mathcal{A}^+(y,p^2) \right. \\
& \left. - \frac{1}{\mathcal{N}} \frac{E_p}{M} \left(1 + \frac{p_z}{M}\right) \mathcal{A}^+_{ON}(y,p^2) \right] g_{1N}\left(\frac{x}{y}\right), \qquad (13)
\end{aligned}
$$

and

$$
\delta^{(N)} g_{1D}(x) \;=\; \frac{1}{8\pi^2 M_D} \int \frac{dy}{y} dp^2 \; \mathcal{A}^+(y,p^2)\; g_{1N}^{off}\left(\frac{x}{y}, p^2\right), \qquad (14)
$$

arise from the off-shell components associated with the deuteron–nucleon vertex and the nucleon structure function, respectively. In the latter $g_{1N}^{off}(x/y,p^2) \equiv \widetilde{g}_{1N}(x/y,p^2) - g_{1N}(x/y)$. The $\delta^{(G)}$ correction is given by:

$$
\begin{aligned}
\delta^{(G)} g_{1D}(x) \;=\; & -\frac{M_D M}{2} \int dy\; dp^2 \; \left[(M_D - 2E_p) \frac{p_z}{M} \left(u - \frac{w}{\sqrt{2}}\right)^2 - \frac{3 M_D p_z}{2M} v_t^2 \right. \\
& \left. + \sqrt{6}\,(M_D - 2E_p) \cos\theta \left(u - \frac{w}{\sqrt{2}}\right) v_t \right] \; p \cdot q \; G_{(p)}(p,q). \qquad (15)
\end{aligned}
$$

To estimate the size of the relativistic corrections requires a model of the nucleon functions $G_{(i)}$. While the scaling behavior of $G_{(i)}$ can be derived from the parton model, their complete evaluation requires model-dependent input for the non-perturbative parton–nucleon physics. For this purpose it is useful to analyze the problem in terms of relativistic quark–nucleon vertex functions, as described in Refs.[13, 14]. We use an ansatz in which a suitable set of spin $S = 0$ and $S = 1$ vertex functions is chosen firstly to parametrize the unpolarized valence nucleon (viz. $u_V + d_V$ and d_V/u_V) and deuteron ($u_V^D + d_V^D$) data. This fixes the normalization and momentum dependence of the vertex functions. The same set is then used to obtain the $S = 0$ and 1 polarized valence distributions Δq_0 and Δq_1. A simple way to relate these to the polarized distributions Δu_V and Δd_V is via the SU(4) spin-flavor symmetric relations [21]: $\Delta u_V = 3\Delta q_0/2 - \Delta q_1/6$ and $\Delta d_V = -\Delta q_1/3$. Of course our formal results do not rely on the use of SU(4) symmetry — these relations merely provide a convenient way to parametrize the polarized quark distributions. Indeed, to reproduce the correct large-x behavior of the proton and neutron structure functions requires that the various spin-flavor quark distributions have different asymptotic x-dependence [22, 23], which necessarily breaks SU(4) symmetry. Within the present approach this is achieved by using different Dirac structures and momentum dependence for the $S = 0$ and $S = 1$ vertices.

We find the polarized and unpolarized data can be fitted with the structures $I\Phi_0^a(p,k)$ and $\not{k}\Phi_0^b(p,k)$ for the scalar vertex, and $\gamma_5\gamma_\alpha\Phi_1(p,k)$ for the pseudovector, where k is the quark momentum. The momentum dependence in the vertices is parametrized by the simple form: $\Phi(p,k) = N(p^2) \cdot k^2/(k^2-\Lambda^2)^n$, with $N(p^2)$ determined through baryon number conservation. A best fit to the experimental nucleon distributions at $Q^2 = 10$ GeV^2 (when evolved from $Q_0^2 \simeq (0.32\text{ GeV})^2$ using leading order evolution[2]) is obtained for cut-off parameters $\Lambda_0^a = 1.0$ GeV and $\Lambda_0^b = 1.1$ GeV, and exponents $n_0^a = 2.0$ and $n_0^b = 2.8$, for the two scalar vertices. These contribute to the total scalar distribution as: $\Delta q_0(x) = r\ \Delta q_0^a(x) + (1-r)\ \Delta q_0^b(x)$, with $r = 0.15$. The parameters for the pseudovector vertex are $\Lambda_1 = 1.8$ GeV and $n_1 = 3.2$. The mass parameters associated with the intermediate spectator states are taken to be $m_{0(1)} = (p-k)^2 = 0.9(1.6)$ GeV.

With these parameters the first moments of the polarized valence distributions in the proton are $\int_0^1 dx\ \Delta u_V(x) = 0.99$ and $\int_0^1 dx\ \Delta d_V(x) = -0.27$, which saturates the Bjorken sum rule: $\int_0^1 dx\ (\Delta u_V(x) - \Delta d_V(x)) = g_A$. The x-dependence of the polarized proton structure function $xg_{1p}(x) = x(4\Delta u_V(x) + \Delta d_V(x))/18$ is shown in Fig.1. In the valence quark dominated region ($x > 0.3$) the result agrees rather well with the SLAC, EMC and SMC proton data [3, 4, 5]. A negatively polarized sea component at $x < 0.3$ would bring the curve even closer to the data points.

For the deuteron, the structure function calculated from Eq.(9) is also shown in Fig.1 (scaled by a factor 1/2). The agreement with the SMC data [1] in the valence region is also quite good. Because the structure function is not very sensitive to explicit p^2-dependence in the quark–nucleon vertex functions, we take $N(p^2)$ to be constant. The numerical values of these normalization constants are fixed through valence quark number conservation in the deuteron to be (1.2, 0.6, 0.8)% smaller (for the Φ_0^a, Φ_0^b, Φ_1 vertices respectively) than in the case of the free nucleon.

The resulting ratio, g_{1D}/g_{1N}, is displayed in Fig.2. For large x it exhibits the same characteristic shape as for the (unpolarized) nuclear EMC effect, namely a dip of $\sim$ 7–8% at $x \approx 0.6$ and a steep rise due to Fermi motion for $x > 0.6$. For small x it stays below unity, where it can be reasonably well approximated by a constant depolarization factor, $1 - 3/2\ \omega_D$, as is typically done in data analyses [1]. Also shown in Fig.2 is the ratio of

[2] While a next-to-leading order analysis is important for a precise determination of the Q^2-dependence of g_1 and the Bjorken sum rule [24], the present treatment is perfectly adequate for the purpose of evaluating the relative sizes of the nuclear corrections.

the convolution ansatz (Eqs.(9) – (11)) to the full calculation (dashed curve). As one can see, this ansatz works remarkably well for most x, the only sizable deviations occurring for $x > 0.8$.

To obtain a better idea about the precise origin of the off-shell effects, we plot in Fig.3 the individual off-shell corrections defined in Eqs.(13)–(15), relative to g_{1D}. For $x < 0.8$ each of the corrections is of order 0.5% or less. As $x \to 1$, however, the convolution-breaking off-shell effects increase by more than an order of magnitude, and will need to be accounted for if one is to obtain accurate information on the $x \to 1$ behavior of the neutron structure function extracted from deuterium data. We have also made the calculation for the case of pseudoscalar πN coupling [19], where we find the total off-shell correction to be roughly twice as large [25]. However pseudoscalar coupling is usually considered to be less realistic than the pseudovector form. The corrections may also depend on the model for the off-shell nucleon structure function, although the fact that the off-shell effects become large at $x \sim 1$ is largely independent of the $x \to 1$ structure of the on-shell nucleon input [25].

Finally, as an estimate of the uncertainty, Δg_{1n}, introduced into the extracted g_{1n} through neglect of various nuclear effects, we plot in Fig.4 the quantity $\Delta g_{1n}(x) \equiv 2 g_{1D}(x) \left(\widetilde{R}(x) - R(x)\right) / \widetilde{R}(x) R(x)$, where $R(x) = g_{1D}(x)/g_{1N}(x)$ is obtained in the full model, while $\widetilde{R}(x)$ is determined through the convolution formula (solid) and via the ansatz $\widetilde{R}(x) = 1 - 3/2\ \omega_D$ (dashed), with $\omega_D = 4.7\%$ [19]. Neglecting off-shell corrections then leads to an uncertainty $|\Delta g_{1n}| \sim 0.002$ in the absolute value of g_{1n}, which is still quite small. For comparison we also illustrate the error associated with using a different value for ω_D in $\widetilde{R}(x) = 1 - 3/2\ \omega_D$, namely 5.8% (dot-dashed) as used in the SMC analysis [1]. A larger value for ω_D would have the effect of shifting the constant curve in Fig.2 down below the "full" result, thereby changing the overall sign of Δg_{1n} in Fig.4. Thus, depending on the precise value of ω_D used, the procedure of applying a constant depolarization factor may lead to an absolute uncertainty $|\Delta g_{1n}| \sim 0.01$, or about 10% of the value of g_{1n} at $x \approx 0.1 - 0.2$.

In summary, we have calculated the polarized deuteron structure function g_{1D} within a covariant framework. Our analysis includes, for the first time, a detailed investigation of effects associated with the off-shell structure of bound nucleons — in addition to the more familiar Fermi motion and binding effects. The conventional convolution model can only be recovered from the full result by taking on-shell limits in the virtual nucleon

structure function and in the nucleon–deuteron interaction. In the end, our conclusion is that the major uncertainty in the extraction of g_{1n} comes from the lack of knowledge of the deuteron D-state wavefunction. The off-shell corrections for $x < 0.7$ are small, of order 0.5% out of a total nuclear effect of $\sim 7\%$. At larger x the correction increases rapidly to $> 5\%$, and will be relevant for higher moment analyses of g_{1n}. Although at present the nuclear effects still lie within the error bars of available data, in the upcoming, high-statistics SLAC E154 and HERMES experiments a careful analysis of all nuclear effects will be necessary.

Acknowledgements.

We would like to thank S.Kulagin and H.Meyer for helpful comments, and W.Weise for a careful reading of the manuscript. W.M. would like to thank the theoretical physics group at the University of Adelaide for its hospitality during a recent visit. This work was supported by the Australian Research Council and the BMFT grant 06 OR 735.

References

[1] SM Collab., B.Adeva et al., Phys.Lett. B 302 (1993) 533; B 320 (1994) 400.

[2] E142 Collab., P.L.Anthony et al., Phys.Rev.Lett. 71 (1993) 959.

[3] G.Baum et al., Phys.Rev.Lett. 51 (1983) 1135.

[4] EM Collab., J.Ashman et al., Phys.Lett. B 206 (1988) 364.

[5] SM Collab., B.Adeva et al., Phys.Lett. B 329 (1994) 399.

[6] E143 Collab., K.Abe et al., preprint SLAC-PUB-6508 (Aug 1994).

[7] H.Khan and P.Hoodbhoy, Phys.Lett. B 298 (1993) 181.

[8] L.L.Frankfurt and M.I.Strikman, Nucl.Phys. A 405 (1983) 557.

[9] R.M.Woloshyn, Nucl.Phys. A 496 (1989) 749.

[10] C.Ciofi degli Atti, E.Pace and G.Salme, Phys.Rev. C 46 (1992) R1591.

[11] M.V.Tokarev, Phys.Lett. B 318 (1993) 559.

[12] L.P.Kaptari et al., Phys.Lett. B 321 (1994) 271; preprint DFUPG 92/94 (1994).

[13] W.Melnitchouk, A.W.Schreiber and A.W.Thomas, Phys.Lett. B 335 (1994) 11.

[14] W.Melnitchouk, A.W.Schreiber and A.W.Thomas, Phys.Rev. D 49 (1994) 1183.

[15] S.A.Kulagin, G.Piller and W.Weise, Phys.Rev. C 50 (1994) 1154.

[16] G.Altarelli and G.G.Ross, Phys.Lett. B 212 (1988) 391; R.D.Carlitz, J.C.Collins and A.H.Mueller, Phys.Lett. B 214 (1988) 229; A.V.Efremov, J.Soffer and O.V.Teryaev, Nucl.Phys. B 346 (1990) 97; R.L.Jaffe and A.Manohar, Nucl.Phys. B 337 (1990) 509; G.M.Shore and G. Veneziano, Nucl.Phys. B 381 (1992) 23; S.D.Bass and A.W.Thomas, Phys.Lett. B 312 (1993) 345; Prog.Part.Nucl.Phys. 33 (1994) 449; S.D.Bass and P.V.Landshoff, Phys.Lett. B 336 (1994) 537.

[17] L.Mankiewicz and A.Schäfer, Phys.Lett. B 274 (1992) 199.

[18] P.Hoodbhoy, R.L.Jaffe and A.Manohar, Nucl.Phys. B 312 (1989) 571.

[19] W.W.Buck and F.Gross, Phys.Rev. D 20 (1979) 2361; R.G.Arnold, C.E.Carlson and F.Gross, Phys.Rev. C 21 (1980) 1426.

[20] J.A.Tjon, Nucl. Phys. A 463 (1987) 157C; D.Plumper and M.F.Gari, Z. Phys. A 343 (1992) 343; F.Gross, J.W.Van Orden and K.Holinde, Phys. Rev. C 45 (1992) 2094.

[21] H.Meyer and P.J.Mulders, Nucl. Phys. A 528 (1991) 589.

[22] F.E.Close and A.W.Thomas, Phys.Lett. B 212 (1988) 227.

[23] S.J.Brodsky, M.Burkardt and I.Schmidt, preprint SLAC-PUB-6087 (1994).

[24] F.E.Close and R.G.Roberts, Phys.Lett. B 316 (1993) 165; J.Ellis and M.Karliner, Phys.Lett. B 313 (1993) 131; G.Altarelli, P.Nason and G.Ridolfi, Phys. Lett. B 320 (1994) 152.

[25] W.Melnitchouk, G.Piller and A.W.Thomas, to be published.

FIGURE CAPTIONS.

1. Valence component of the proton (solid) and deuteron (dashed) g_1 structure functions at $Q^2 = 10$ GeV^2. The data are from Refs.[1, 3, 4, 5, 6]. The deuteron data are scaled by a factor 1/2.

2. Ratio of deuteron and nucleon structure functions in the full model (solid), and with a constant depolarization factor $1 - 3/2\ \omega_D$ (dotted), with $\omega_D = 4.7\%$ [19]. Dashed curve is ratio of g_{1D} calculated via convolution and in the full model.

3. Nucleon off-shell corrections to g_{1D}: $\delta^{(N)} g_{1D}$ (dotted), $\delta^{(\mathcal{A})} g_{1D}$ (dashed), $\delta^{(G)} g_{1D}$ (dot-dashed) and the sum (solid), as a fraction of the total g_{1D}.

4. Estimate of the error, Δg_{1n} (scaled by a factor 10), introduced into g_{1n} extracted from g_{1D} by neglecting off-shell effects (solid), and by using a constant depolarization factor $(1 - 3/2\ \omega_D)$, with $\omega_D = 4.7\%$ [19] (dashed) and 5.8% [1] (dot-dashed). To show the significance of the correction, $\Delta g_{1n}(\times 10)$ is compared with the SLAC E142 g_{1n} data [2].

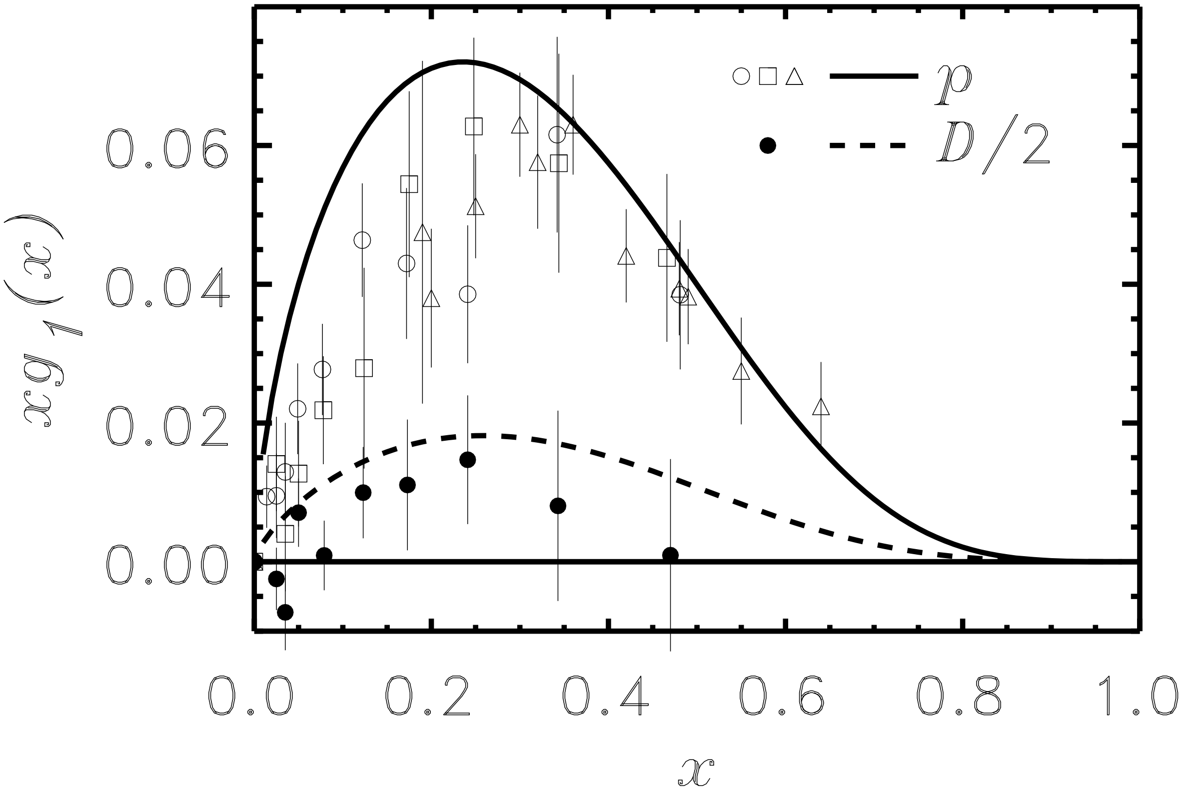
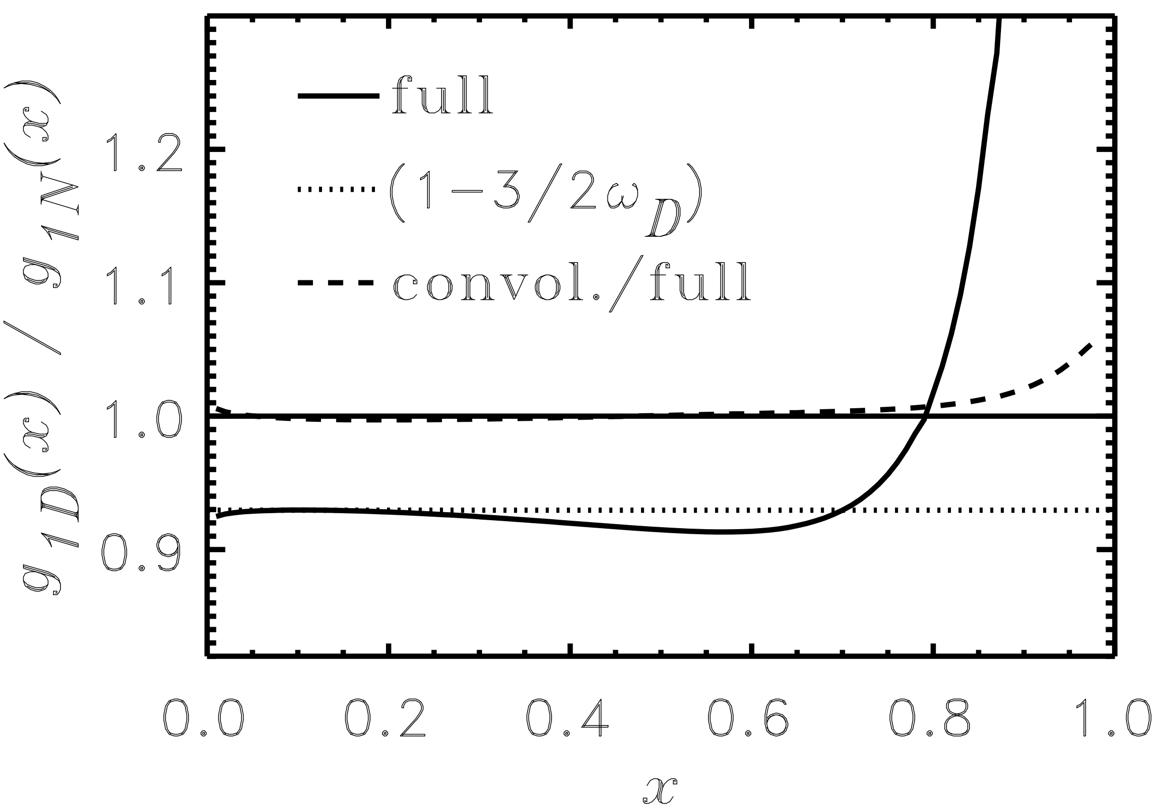


Fig.1

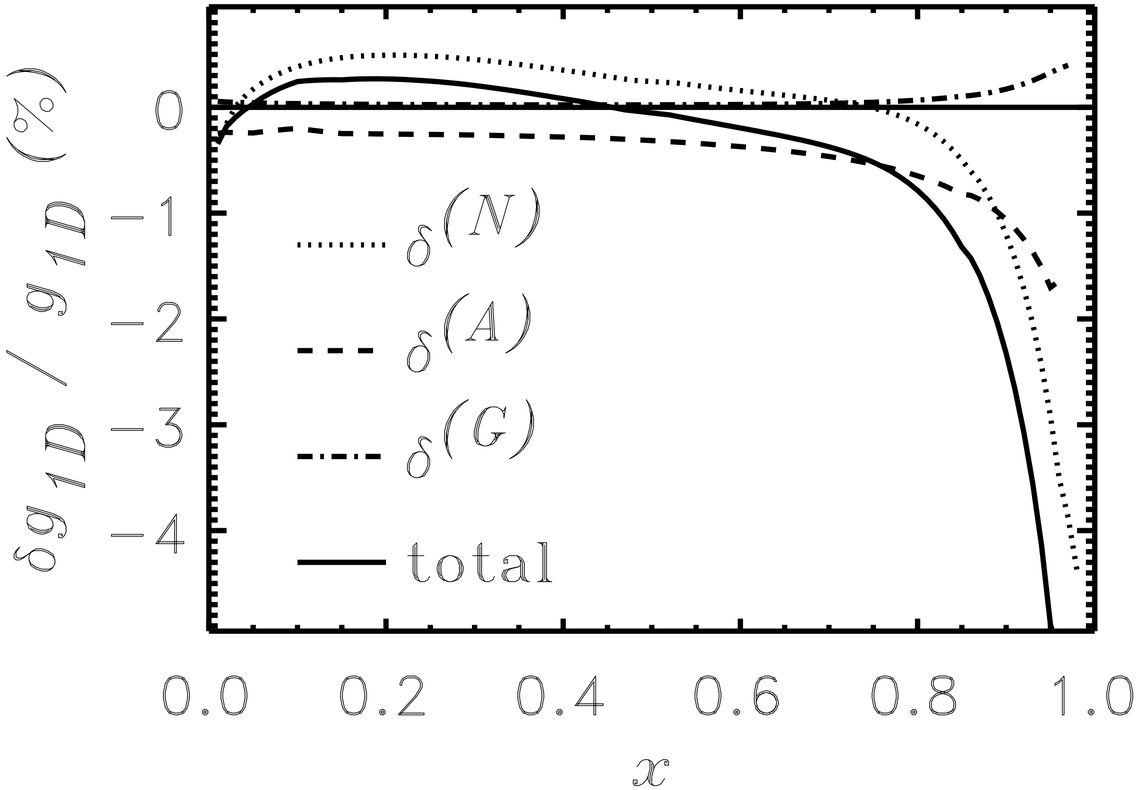


full
$(1-3/2\omega_D)$
convol./full
1.2
1.1
1.0
0.9
$g_{1D}(x)$ / $g_{1N}(x)$
0.0
0.2
0.4
0.6
0.8
1.0
x

Fig.2

This figure "fig1-2.png" is available in "png" format from:

http://arXiv.org/ps/hep-ph/9501282v1

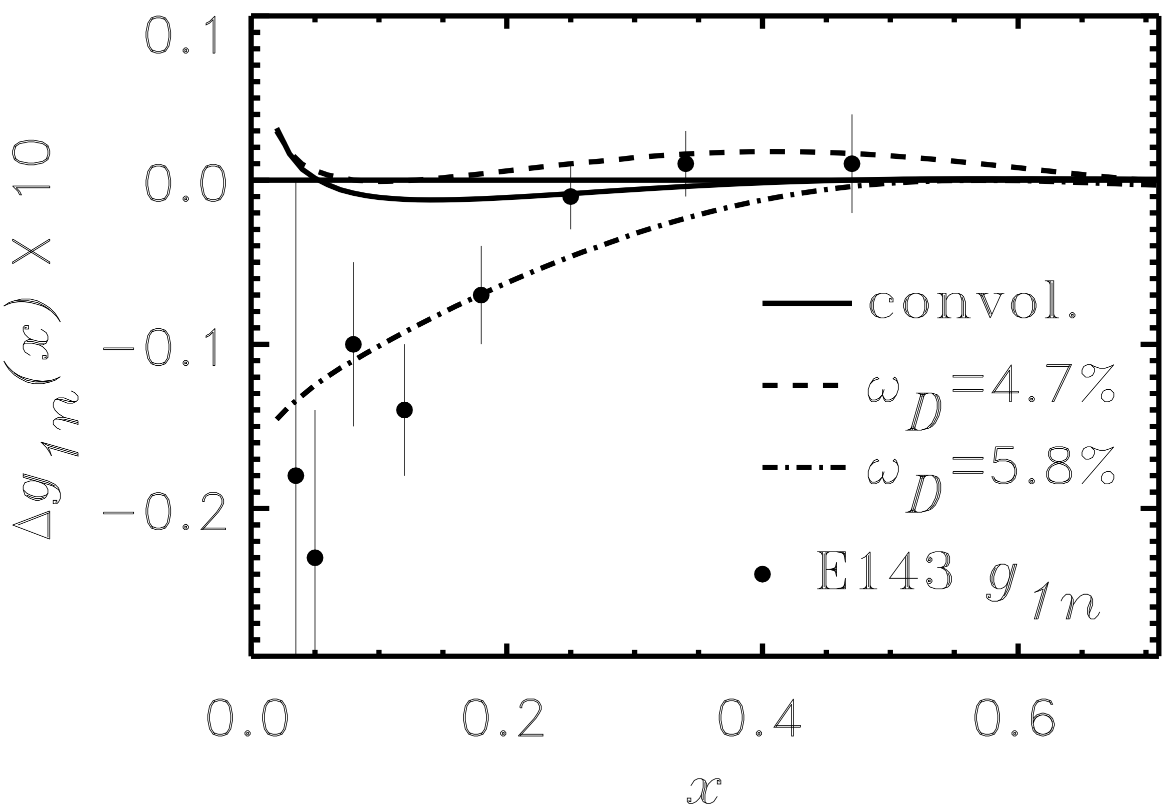


δg_{1D} / g_{1D} (%)
0
−1
−2
−3
−4
$\delta^{(N)}$
$\delta^{(A)}$
$\delta^{(G)}$
total
0.0
0.2
0.4
0.6
0.8
1.0
x

Fig.3

This figure "fig1-3.png" is available in "png" format from:

http://arXiv.org/ps/hep-ph/9501282v1



0.1
0.0
−0.1
−0.2
$\Delta g_{1n}(x)$ X 10
convol.
ω_D=4.7%
ω_D=5.8%
E143 g_{1n}
0.0
0.2
0.4
0.6
x

Fig.4